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Finite Theories after the Discovery of a Higgs-like Boson at the LHC

Sven Heinemeyer[1] *, Myriam Mondragón[2] †
and George Zoupanos[3] ‡§

[1]*Instituto de Física de Cantabria (CSIC-UC)*
Edificio Juan Jorda, Avda. de Los Castros s/n
39005 Santander, Spain

[2] *Instituto de Física*
Universidad Nacional Autónoma de México
Apdo. Postal 20-364, México 01000 D.F., México

[3] *Max-Planck-Institut für Physik (Werner-Heisenberg-Institut)*
Föhringer Ring 6 80805 München and
Arnold-Sommerfeld-Center für Theoretische Physik
Department für Physik, Ludwig-Maximilians-Universität München
Theresienstrasse 37, 80333 Muenchen, Germany and
Theory Group, Physics Department
CERN, Geneva, Switzerland

Abstract

Finite Unified Theories (FUTs) are $N = 1$ supersymmetric Grand Unified Theories (GUTs) which can be made finite to all-loop orders, based on the principle of reduction of couplings, and therefore are provided with a large predictive power. Confronting the predictions of SU(5) FUTs with the top and bottom quark masses and other low-energy experimental constraints a light Higgs-boson mass in the range $M_h \sim 121 - 126$ GeV was predicted, in striking agreement with the recent discovery of a Higgs-like state around ~ 125.7 GeV at ATLAS and CMS. Furthermore the favoured model, a finiteness constrained version of the MSSM, naturally predicts a relatively heavy spectrum

* email: Sven.Heinemeyer@cern.ch
† email: myriam@fisica.unam.mx
‡ email: George.Zoupanos@cern.ch
§On leave from Physics Department, National Technical University, Zografou Campus: Heroon Polytechniou 9, 15780 Zografou, Athens, Greece

with coloured supersymmetric particles above ~ 1.5 TeV, consistent with the non-observation of those particles at the LHC. Restricting further the best FUT's parameter space according to the discovery of a Higgs-like state and B-physics observables we find predictions for the rest of the Higgs masses and the s-spectrum.

1 Introduction

The success of the Standard Model (SM) of Elementary Particle Physics has recently been confirmed by the the observation of a state compatible with a (SM-like) Higgs boson at the LHC [1]. Still, the number of free parameters of the SM points towards the possibility that it is the low energy limit of a more fundamental theory. One of the most studied extensions of the SM is the Minimal Supersymmteric Standard Model (MSSM) [2], where one particular realization is the constrained MSSM (CMSSM) [3] with only five free parameters. Recent LHC results discard some regions of the CMSSM and point towards a heavy spectrum in case this particular version of SUSY is realized in nature [4].

Searching for renormalization group invariant (RGI) relations [5–16] holding below the Planck scale down to the GUT scale provides a different strategy to search for a more fundamental theory, whose basic ingredients are GUTs and supersymmetry (SUSY), and with far reaching consequences [6–9]. An outstanding feature of the use of RGI relations is that one can guarantee their validity to all-orders in perturbation theory by studying the uniqueness of the resulting relations at one-loop [10]. Even more remarkable is the fact that it is possible to find RGI relations among couplings that guarantee finiteness to all-orders in perturbation theory [11].

The Gauge–Yukawa unification scheme, based in RGI relations applied in the dimensionless couplings of supersymmetric GUTs, such as gauge and Yukawa couplings, had noticeable successes by predicting correctly the top quark mass in the finite [6] and in the minimal $N = 1$ supersymmetric SU(5) GUTs [7]. Finite Unified Theories are $N = 1$ supersymmetric GUTs which can be made finite to all-loop orders, including the soft-SUSY breaking sector (for reviews and detailed refs. see [9, 12–15]), which involves parameters of dimension one and two. Taking into account the restrictions resulting from the low-energy observables, it was possible to extend the predictive power of the RGI method to the Higgs sector and the SUSY spectrum. the Higgs

boson mass thus eventually predicted [16]

$$M_h \simeq 121 - 126 \text{ GeV} \tag{1}$$

is in agreement with the recent discovery Higgs-like state at the LHC [1]. As further features a heavy SUSY spectrum and large values of $\tan\beta$ (the ratio of the two vacuum expectation values) were found [16].

In this letter, first we review two $SU(5)$-based finite SUSY models and their predictions, taking into account the restrictions resulting from the low-energy observables [16]. Only one model survives all the phenomenological constraints. Then we extend our previous analysis by imposing more recent contraints resulting from the bounds on $\text{BR}(B_s \to \mu^+\mu^-)$. Moreover, as the crucial new ingredient we interpret the newly discoverd particle at ~ 126 GeV as the lightest MSSM Higgs boson and we analyse which constraints imposes the measured value of the Higgs boson mass on the predictions of the SUSY spectrum.

2 Finiteness

Finiteness can be understood by considering a chiral, anomaly free, $N = 1$ globally supersymmetric gauge theory based on a group G with gauge coupling constant g. The superpotential of the theory is given by

$$W = \frac{1}{2}\, m^{ij}\, \Phi_i\, \Phi_j + \frac{1}{6}\, C^{ijk}\, \Phi_i\, \Phi_j\, \Phi_k\ , \tag{2}$$

where m^{ij} (the mass terms) and C^{ijk} (the Yukawa couplings) are gauge invariant tensors and the matter field Φ_i transforms according to the irreducible representation R_i of the gauge group G. All the one-loop β-functions of the theory vanish if the β-function of the gauge coupling $\beta_g^{(1)}$, and the anomalous dimensions of the Yukawa couplings $\gamma_i^{j(1)}$, vanish, i.e.

$$\sum_i \ell(R_i) = 3C_2(G)\,,\ \frac{1}{2}C_{ipq}C^{jpq} = 2\delta_i^j g^2 C_2(R_i)\ , \tag{3}$$

where $\ell(R_i)$ is the Dynkin index of R_i, and $C_2(G)$ is the quadratic Casimir invariant of the adjoint representation of G. These conditions are also enough to guarantee two-loop finiteness [17]. A striking fact is the existence of a theorem [11] that guarantees the vanishing of the β-functions to all-orders in

perturbation theory. This requires that, in addition to the one-loop finiteness conditions (3), the Yukawa couplings are reduced in favour of the gauge coupling to all-orders (see [15] for details). Alternatively, similar results can be obtained [18] using an analysis of the all-loop NSVZ gauge beta-function [19].

Next consider the superpotential given by (2) along with the Lagrangian for soft supersymmetry breaking (SSB) terms

$$-\mathcal{L}_{\rm SB} = \frac{1}{6}\,h^{ijk}\,\phi_i\phi_j\phi_k + \frac{1}{2}\,b^{ij}\,\phi_i\phi_j + \frac{1}{2}\,(m^2)^j_i\,\phi^{*\,i}\phi_j + \frac{1}{2}\,M\,\lambda\lambda + \text{h.c.}, \quad (4)$$

where the ϕ_i are the scalar parts of the chiral superfields Φ_i, λ are the gauginos and M their unified mass, h^{ijk} and b^{ij} are the trilinear and bilinear dimensionful couplings respectively, and $(m^2)^j_i$ the soft scalars masses. Since we would like to consider only finite theories here, we assume that the gauge group is a simple group and the one-loop β-function of the gauge coupling g vanishes. We also assume that the reduction equations admit power series solutions of the form

$$C^{ijk} = g\sum_n \rho^{ijk}_{(n)} g^{2n}\ . \quad (5)$$

According to the finiteness theorem of ref. [11,20], the theory is then finite to all orders in perturbation theory, if, among others, the one-loop anomalous dimensions $\gamma_i^{j(1)}$ vanish. the one- and two-loop finiteness for h^{ijk} can be achieved through the relation [21]

$$h^{ijk} = -MC^{ijk} + \cdots = -M\rho^{ijk}_{(0)}\,g + O(g^5)\ , \quad (6)$$

where ... stand for higher order terms.

In addition it was found that the RGI SSB scalar masses in Gauge-Yukawa unified models satisfy a universal sum rule at one-loop [22]. This result was generalized to two-loops for finite theories [14], and then to all-loops for general Gauge-Yukawa and finite unified theories [23]. From these latter results, the following soft scalar-mass sum rule is found [14]

$$\frac{(\ m_i^2 + m_j^2 + m_k^2\ )}{MM^\dagger} = 1 + \frac{g^2}{16\pi^2}\,\Delta^{(2)} + O(g^4) \quad (7)$$

for i, j, k with $\rho^{ijk}_{(0)} \neq 0$, where $m^2_{i,j,k}$ are the scalar masses and $\Delta^{(2)}$ is the two-loop correction

$$\Delta^{(2)} = -2\sum_l [(m_l^2/MM^\dagger) - (1/3)]\ \ell(R_l), \quad (8)$$

which vanishes for the universal choice, i.e. when all the soft scalar masses are the same at the unification point. This correction also vanishes in the models considered here.

3 $SU(5)$ Finite Unified Theories

Finite Unified Models have been studied for already two decades. A realistic two-loop finite $SU(5)$ model was presented in [24], and shortly afterwards the conditions for finiteness in the soft susy breaking sector at one-loop [17] were given. Since finite models usually have an extended Higgs sector, in order to make them viable a rotation of the Higgs sector was proposed [25]. the first all-loop finite theory was studied in [6], without taking into account the soft breaking terms. Naturally, the concept of finiteness was extended to the soft breaking sector, where also one-loop finiteness implies two-loop finiteness [21], and then finiteness to all-loops in the soft sector of realistic models was studied [26, 27], although the universality of the soft breaking terms lead to a charged lightest SUSY particle (LSP). This fact was also noticed in [28], where the inclusion of an extra parameter in the Higgs sector was introduced to alleviate it. With the derivation of the sum-rule in the soft supersymmetry breaking sector and the proof that it can be made all-loop finite the construction of all-loop phenomenologically viable finite models was made possible [14, 23].

Here we will examine two all-loop Finite Unified theories with $SU(5)$ gauge group, where the reduction of couplings has been applied to the third generation of quarks and leptons. An extension to three families, and the generation of quark mixing angles and masses in Finite Unified Theories has been addressed in [29], where several examples are given. These extensions are not considered here. Realistic Finite Unified Theories based on product gauge groups, where the finiteness implies three generations of matter, have also been studied [30].

The particle content of the models we will study consists of the following supermultiplets: three $(\overline{\mathbf{5}} + \mathbf{10})$, needed for each of the three generations of quarks and leptons, four $(\overline{\mathbf{5}} + \mathbf{5})$ and one $\mathbf{24}$ considered as Higgs supermultiplets. When the gauge group of the finite GUT is broken the theory is no longer finite, and we will assume that we are left with the MSSM.

Thus, a predictive Gauge-Yukawa unified $SU(5)$ model which is finite to all orders, in addition to the requirements mentioned already, should also

have the following properties:

1. One-loop anomalous dimensions are diagonal, i.e., $\gamma_i^{(1)\,j} \propto \delta^j_i$.
2. Three fermion generations, in the irreducible representations $\overline{\mathbf{5}}_i, \mathbf{10}_i$ ($i = 1,2,3$), which obviously should not couple to the adjoint $\mathbf{24}$.
3. The two Higgs doublets of the MSSM should mostly be made out of a pair of Higgs quintet and anti-quintet, which couple to the third generation.

The two versions of the all-order finite model we will discuss here are the following: the model of [6], which will be labelled **A**, and a slight variation of this model (labelled **B**), which can also be obtained from the class of the models suggested in [26] with a modification to suppress non-diagonal anomalous dimensions.

The superpotential which describes the two models, which we will label **A** and **B**, takes the form [6,14]

$$\begin{aligned} W &= \sum_{i=1}^{3} \Big[\frac{1}{2} g_i^u \,\mathbf{10}_i \mathbf{10}_i H_i + g_i^d \,\mathbf{10}_i \overline{\mathbf{5}}_i \,\overline{H}_i \,\Big] \\ &+ g_{23}^u \,\mathbf{10}_2 \mathbf{10}_3 H_4 + g_{23}^d \,\mathbf{10}_2 \overline{\mathbf{5}}_3 \,\overline{H}_4 + g_{32}^d \,\mathbf{10}_3 \overline{\mathbf{5}}_2 \,\overline{H}_4 \\ &+ \sum_{a=1}^{4} g_a^f \, H_a \,\mathbf{24}\, \overline{H}_a + \frac{g^\lambda}{3} \,(\mathbf{24})^3 \ , \end{aligned} \tag{9}$$

where H_a and $\overline{H}_a$ $(a = 1, \ldots, 4)$ stand for the Higgs quintets and anti-quintets.

The main difference between model **A** and model **B** is that two pairs of Higgs quintets and anti-quintets couple to the $\mathbf{24}$ in **B**, so that it is not necessary to mix them with H_4 and $\overline{H}_4$ in order to achieve the triplet-doublet splitting after the symmetry breaking of $SU(5)$ [14]. Thus, although the particle content is the same, the solutions to the finiteness equations and the sum rules are different, which has repercussions in the phenomenology.

FUTA

After the reduction of couplings the symmetry of the superpotential W (9) is enhanced (for details see [31]). the superpotential for this model is

$$W = \sum_{i=1}^{3} [\ \frac{1}{2} g_i^u\, \mathbf{10}_i \mathbf{10}_i H_i + g_i^d\, \mathbf{10}_i \overline{\mathbf{5}}_i\, \overline{H}_i\ ] + g_4^f\, H_4\, \mathbf{24}\, \overline{H}_4 + \frac{g^\lambda}{3}\, (\mathbf{24})^3 \tag{10}$$

The non-degenerate and isolated solutions to $\gamma_i^{(1)} = 0$ for model **FUTA**, which are the boundary conditions for the Yukawa couplings at the GUT scale, are:

$$\begin{aligned} &(g_1^u)^2 = \frac{8}{5}\, g^2\ ,\ (g_1^d)^2 = \frac{6}{5}\, g^2\ ,\ (g_2^u)^2 = (g_3^u)^2 = \frac{8}{5}\, g^2\ , \\ &(g_2^d)^2 = (g_3^d)^2 = \frac{6}{5}\, g^2\ ,\ (g_{23}^u)^2 = 0\ ,\ (g_{23}^d)^2 = (g_{32}^d)^2 = 0\ , \\ &(g^\lambda)^2 = \frac{15}{7} g^2\ ,\ (g_2^f)^2 = (g_3^f)^2 = 0\ ,\ (g_1^f)^2 = 0\ ,\ (g_4^f)^2 = g^2\ . \end{aligned} \tag{11}$$

In the dimensionful sector, the sum rule gives us the following boundary conditions at the GUT scale for this model [14]:

$$m_{H_u}^2 + 2 m_{\mathbf{10}}^2 = m_{H_d}^2 + m_{\overline{\mathbf{5}}}^2 + m_{\mathbf{10}}^2 = M^2\ , \tag{12}$$

and thus we are left with only three free parameters, namely $m_{\overline{\mathbf{5}}} \equiv m_{\overline{\mathbf{5}}_3}$, $m_{\mathbf{10}} \equiv m_{\mathbf{10}_3}$ and M.

FUTB

Also in the case of **FUTB** the symmetry is enhanced after the reduction of couplings, with the following superpotential [31]

$$\begin{aligned} W = & \sum_{i=1}^{3} [\ \frac{1}{2} g_i^u\, \mathbf{10}_i \mathbf{10}_i H_i + g_i^d\, \mathbf{10}_i \overline{\mathbf{5}}_i\, \overline{H}_i\ ] + g_{23}^u\, \mathbf{10}_2 \mathbf{10}_3 H_4 \\ & + g_{23}^d\, \mathbf{10}_2 \overline{\mathbf{5}}_3\, \overline{H}_4 + g_{32}^d\, \mathbf{10}_3 \overline{\mathbf{5}}_2\, \overline{H}_4 + g_2^f\, H_2\, \mathbf{24}\, \overline{H}_2 + g_3^f\, H_3\, \mathbf{24}\, \overline{H}_3 + \frac{g^\lambda}{3}\, (\mathbf{24})^3\ , \end{aligned} \tag{13}$$

For this model the non-degenerate and isolated solutions to $\gamma_i^{(1)} = 0$ give us:

$$
\begin{aligned}
&(g_1^u)^2 = \frac{8}{5}\, g^2\ ,\ (g_1^d)^2 = \frac{6}{5}\, g^2\ ,\ (g_2^u)^2 = (g_3^u)^2 = \frac{4}{5}\, g^2\ , \qquad (14)\\
&(g_2^d)^2 = (g_3^d)^2 = \frac{3}{5}\, g^2\ ,\ (g_{23}^u)^2 = \frac{4}{5}\, g^2\ ,\ (g_{23}^d)^2 = (g_{32}^d)^2 = \frac{3}{5}\, g^2\ ,\\
&(g^\lambda)^2 = \frac{15}{7} g^2\ ,\ (g_2^f)^2 = (g_3^f)^2 = \frac{1}{2}\, g^2\ ,\ (g_1^f)^2 = 0\ ,\ (g_4^f)^2 = 0\ ,
\end{aligned}
$$

and from the sum rule we obtain:

$$
\begin{aligned}
m_{H_u}^2 + 2m_{\mathbf{10}}^2 &= M^2\ ,\ m_{H_d}^2 - 2m_{\mathbf{10}}^2 = -\frac{M^2}{3}\ ,\\
m_{\overline{\mathbf{5}}}^2 + 3m_{\mathbf{10}}^2 &= \frac{4M^2}{3}\ , \qquad (15)
\end{aligned}
$$

i.e., in this case we have only two free parameters $m_{\mathbf{10}} \equiv m_{\mathbf{10}_3}$ and M for the dimensionful sector.

As already mentioned, after the $SU(5)$ gauge symmetry breaking we assume we have the MSSM, i.e. only two Higgs doublets. This can be achieved by introducing appropriate mass terms that allow to perform a rotation of the Higgs sector [6, 24, 25, 32], in such a way that only one pair of Higgs doublets, coupled mostly to the third family, remains light and acquire vacuum expectation values. To avoid fast proton decay the usual fine tuning to achieve doublet-triplet splitting is performed. Notice that, although similar, the mechanism is not identical to minimal $SU(5)$, since we have an extended Higgs sector.

Thus, after the gauge symmetry of the GUT theory is broken we are left with the MSSM, with the boundary conditions for the third family given by the finiteness conditions, while the other two families are not restricted.

We will now examine the phenomenology of such all-loop Finite Unified theories with $SU(5)$ gauge group and, for the reasons expressed above, we will concentrate only on the third generation of quarks and leptons.

4 Predictions of Low Energy Parameters

Since the gauge symmetry is spontaneously broken below $M_{\rm GUT}$, the finiteness conditions do not restrict the renormalization properties at low energies, and all it remains are boundary conditions on the gauge and Yukawa couplings (11) or (14), the $h = -MC$ (6) relation, and the soft scalar-mass

sum rule at $M_{\rm GUT}$, as applied in the two models, Eq. (12) or (15). Thus we examine the evolution of these parameters according to their RGEs up to two-loops for dimensionless parameters and at one-loop for dimensionful ones with the relevant boundary conditions. Below $M_{\rm GUT}$ their evolution is assumed to be governed by the MSSM. We further assume a unique supersymmetry breaking scale M_s (which we define as the geometric mean of the stop masses) and therefore below that scale the effective theory is just the SM.

We now briefly review the comparison of the predictions of the two models (**FUTA**, **FUTB**) with the experimental data, starting with the heavy quark masses see ref. [16] for more details.

We use for the top quark the value for the pole mass [33]

$$m_t^{\rm exp} = (173.2 \pm 0.9)\ \text{GeV}\ , \tag{16}$$

and we recall that the theoretical prediction for m_t of the present framework may suffer from a correction of $\sim 4\%$ [9, 12, 34, 35]. For the bottom quark mass we use the value at M_Z [36]

$$m_b(M_Z) = (2.83 \pm 0.10)\ \text{GeV}, \tag{17}$$

to avoid uncertainties that come from the further running from the M_Z to the m_b mass.

In fig.1 we show the **FUTA** and **FUTB** predictions for m_t and $m_b(M_Z)$ as a function of the unified gaugino mass M, for the two cases $\mu < 0$ and $\mu > 0$. In the value of the bottom mass m_b, we have included the corrections coming from bottom squark-gluino loops and top squark-chargino loops [37], known usually as the Δ_b effects. The bounds on the $m_b(M_Z)$ and the m_t mass clearly single out **FUTB** with $\mu < 0$, as the solution most compatible with this experimental constraints. Although $\mu < 0$ is already challenged by present data of the anomalous magnetic moment of the muon a_μ [38, 39], a heavy SUSY spectrum as the one we have here (see below) gives results for a_μ very close to the SM result, and thus cannot be excluded.

We now analyze the impact of further low-energy observables on the model **FUTB** with $\mu < 0$. As additional constraints we consider the following observables: the rare b decays $\text{BR}(b \to s\gamma)$ and $\text{BR}(B_s \to \mu^+\mu^-)$.

For the branching ratio $\text{BR}(b \to s\gamma)$, we take the value given by the Heavy Flavour Averaging Group (HFAG) is [40]

$$\text{BR}(b \to s\gamma) = (3.55 \pm 0.24^{+0.09}_{-0.10} \pm 0.03) \times 10^{-4}. \tag{18}$$

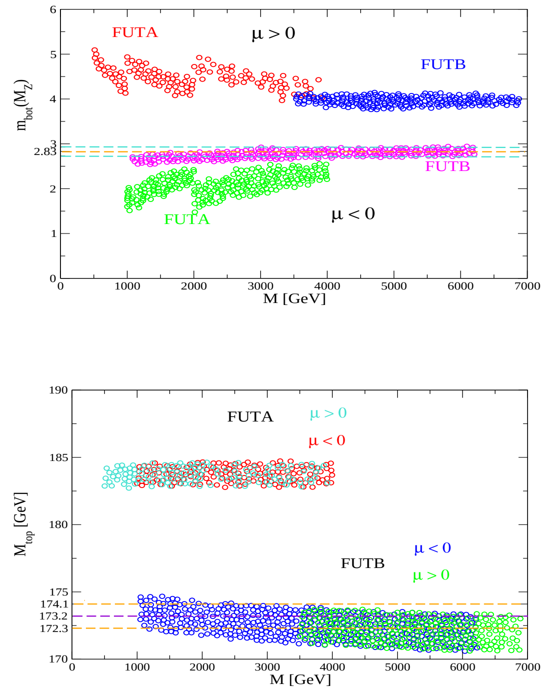

Figure 1: The bottom quark mass at the Z boson scale (upper) and top quark pole mass (lower plot) are shown as function of M, the unified gaugino mass, for both models.

For the branching ratio $\mathrm{BR}(B_s \to \mu^+\mu^-)$, the SM prediction is at the level of 10^{-9}, while the present experimental upper limit is

$$\mathrm{BR}(B_s \to \mu^+\mu^-) = 4.5 \times 10^{-9} \tag{19}$$

at the 95% C.L. [41]. ¶

For the lightest Higgs mass prediction we use the code `FeynHiggs` [43–45]. The prediction for M_h of **FUTB** with $\mu < 0$ is shown in Fig. 2, where the

¶While we were finalizing this paper, a first measurement at the $\sim 3\sigma$ level of $\mathrm{BR}(B_s \to \mu^+\mu^-)$ was published by the LHCb collaboration [42]. the value is given as $\mathrm{BR}(B_s \to \mu^+\mu^-) = (3.2^{+1.4}_{-1.2}(\mathrm{stat})^{+0.5}_{-0.3}(\mathrm{syst})) \times 10^{-9}$, i.e. the upper limit at the 95% CL is slightly higher than what we used as an upper limit. Furthermore, no combination of this new result with the existing limits exists yet. Consequently, as we do not expect a sizable impact of the very new measurement on our results, we stick for this analysis to the simple upper limit.

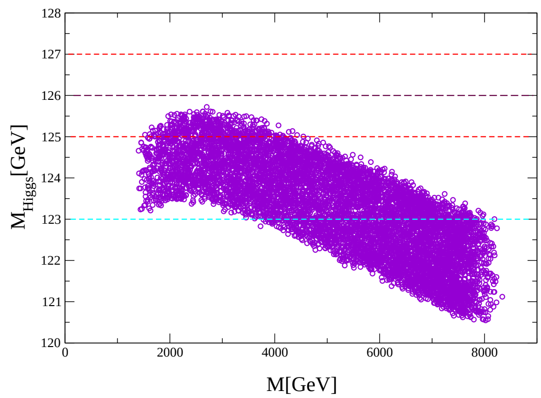

Figure 2: The lightest Higgs mass, M_h, as function of M for the model **FUTB** with $\mu < 0$, see text.

constraints from the two B physics observables are taken into account. The lightest Higgs mass ranges in

$$M_h \sim 121 - 126 \text{ GeV} , \tag{20}$$

where the uncertainty comes from variations of the soft scalar masses. To this value one has to add at least ± 2 GeV coming from unkonwn higher order corrections [44]. We have also included a small variation, due to threshold corrections at the GUT scale, of up to 5% of the FUT boundary conditions. the masses of the heavier Higgs bosons are found at higher values in comparison with our previous analyses [16, 46]. This is due to the more stringent bound on BR($B_s \to \mu^+\mu^-$), which pushes the heavy Higgs masses beyond ~ 1 TeV, excluding their discovery at the LHC. We furthermore find in our analysis that the lightest observable SUSY particle (LOSP) is either the stau or the second lightest neutralino, with mass starting around ~ 500 GeV.

As the crucial new ingredient we take into account the recent observations of a Higgs-like state discovered at LHC. We impose a constraint on our results to the Higgs mass of

$$M_h \sim 126.0 \pm 1 \pm 2 \text{ GeV} , \tag{21}$$

where ± 1 comes from the experimental error and ± 2 corresponds to the theoretical error, and see how this affects the SUSY spectrum. Constraining the allowed values of the Higgs mass this way puts a limit on the allowed values of the unified gaugino mass, as can be seen from Fig. 2. the red lines correspond

to the pure experimental uncertainty and restrict 2 TeV $\lesssim M \lesssim$ 5 TeV. the blue line includes the additional theory uncertainty of ± 2 GeV. Taking this uncertainty into account no bound on M can be placed. However, a substantial part of the formerly allowed parameter points are now excluded. This in turn restricts the lightest observable SUSY particle (LOSP), which turns out to be the light scalar tau. In Fig. 3 the effects on the mass of the LOSP are demonstrated. Without any Higgs mass constraint all coloured points are allowed. Imposing $M_h = 126 \pm 1$ GeV only the green (light shaded) points are allowed, restricting the mass to be between about 500 GeV and 2500 GeV. the lower values might be experimentally accessible at the ILC with 1000 GeV center-of-mass energy or at CLIC with an energy up to ~ 3 TeV. Taking into account the theory uncertainty on M_h also the blue (dark shaded) points are allowed, permitting the LOSP mass up to ~ 4 TeV. If the upper end of the parameter space were realized the light scalar tau would remain unobservable even at CLIC.

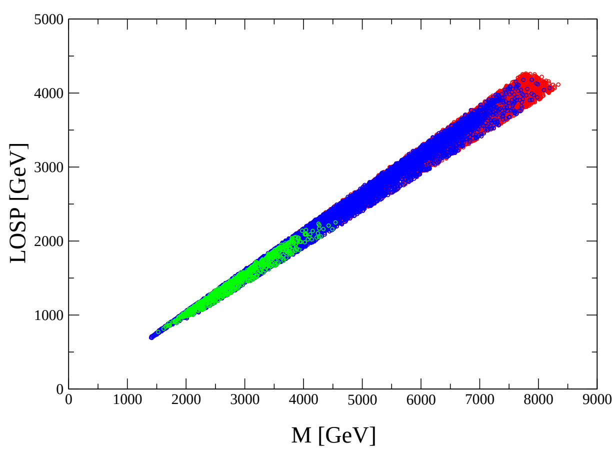

Figure 3: The mass of the LOSP is presented as a function of M. Shown are only points that fulfill the B physics constraints. The green (light shaded) points correspond to $M_h = 126 \pm 1$ GeV, the blue (dark shaded) points have $M_h = 126 \pm 3$ GeV, and the red points have no M_h restriction.

The full particle spectrum of model **FUTB** with $\mu < 0$, compliant with quark mass constraints and the B-physics observables is shown in Fig. 4. In the upper (lower) plot we impose $M_h = 126 \pm 3(1)$ GeV. Without any

Mbot(M_Z)	2.74	Mtop	174.1
Mh	125.0	MA	1517
MH	1515	MH$^{\pm}$	1518
Stop1	2483	Stop2	2808
Sbot1	2403	Sbot2	2786
Mstau1	892	Mstau2	1089
Char1	1453	Char2	2127
Neu1	790	Neu2	1453
Neu3	2123	Neu4	2127
Mgluino	3632		

Table 1: A representative spectrum of a light **FUTB**, $\mu < 0$ spectrum, compliant with the B physics constraints. All masses are in GeV.

M_h restrictions the coloured SUSY particles have masses above ~ 1.8 TeV in agreement with the non-observation of those particles at the LHC [47]. Including the Higgs mass constraints in general favours the lower part of the SUSY particle mass spectra, but also cuts away the very low values. Neglecting the theory uncertainties of M_h (as shown in the lower plot of Fig. 4) permits SUSY masses which would remain unobservable at the LHC, the ILC or CLIC. On the other hand, large parts of the allowed spectrum of the lighter scalar tau or the lighter neutralinos might be accessible at CLIC with $\sqrt{s} = 3$ TeV. Including the theory uncertainties, even higher masses are permitted, further weakening the discovery potential of the LHC and future e^+e^- colliders. A numerical example of the lighter part of the spectrum is shown in Table 1. If such a spectrum were realized, the coloured particles are at the border of the discovery region at the LHC. Some uncoloured particles like the scalar taus, the light chargino or the lighter neutralinos would be in the reach of a high-energy Linear Collider.

5 Conclusions

We examined the predictions of two $SU(5)$ Finite Unified Theories in light of the recent discovery of a Higgs-like state at the LHC and on the new bounds on the branching ratio $\mathrm{BR}(B_s \to \mu^+\mu^-)$. Only one model is consistent with all the phenomenological constraints. Compared to our previous analysis [16],

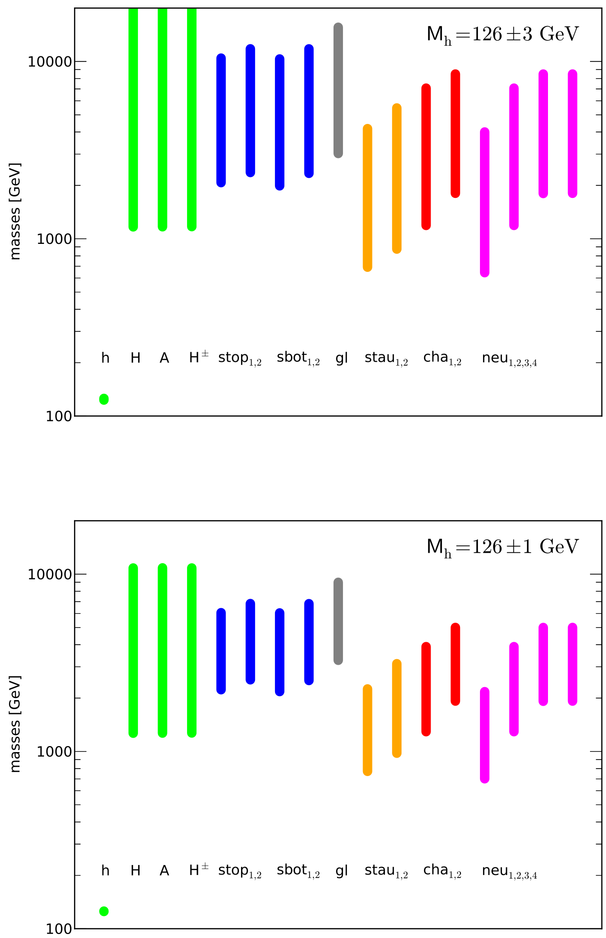

Figure 4: The upper (lower) plot shows the spectrum after imposing the constraint $M_h = 126 \pm 3\,(1)$ GeV. The particle spectrum of model **FUTB** with $\mu < 0$, where the points shown are in agreement with the quark mass constraints and the B-physics observables. The light (green) points on the left are the various Higgs boson masses. The dark (blue) points following are the two scalar top and bottom masses, followed by the lighter (gray) gluino mass. Next come the lighter (beige) scalar tau masses. The darker (red) points to the right are the two chargino masses followed by the lighter shaded (pink) points indicating the neutralino masses.

the new bound on $\mathrm{BR}(B_s \to \mu^+\mu^-)$ excludes values for the heavy Higgs bosons masses below 1 $\sim$ TeV, and in general allows only a very heavy SUSY spectrum. The Higgs mass constraint favours the lower part of this spectrum, with SUSY masses ranging from $\sim$ 500 GeV up to the multi-TeV level, where the lower part of the spectrum could be accessible at the ILC or CLIC.

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